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Hadronic Weak Decays of Heavy Mesons and Nonfactorization

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Abstract

The parameters $\chi_{1,2}$, which measure nonfactorizable soft gluon contributions to hadronic weak decays of mesons, are updated by extracting them from the data of $D,\,B\to PP,\ VP$ decays (P: pseudoscalar meson, V: vector meson). It is found that χ_2 ranges from -0.36 to -0.60 in the decays from $D\to\bar{K}\pi$ to $D^+\to\phi\pi^+,\ D\to\bar{K}^*\pi$, while it is of order 10% with a positive sign in $B\to\psi K,\ D\pi,\ D^*\pi,\ D\rho$ decays. Therefore, the effective parameter a_2 is process dependent in charm decay, whereas it stays fairly stable in B decay. This implies the picture that nonfactorizable effects become stronger when the decay particles become less energetic after hadronization. As for $D,\,B\to VV$ decays, the presence of nonfactorizable terms in general prevents a possible definition of effective a_1 and a_2. This is reenforced by the observation of a large longitudinal polarization fraction in $B\to\psi K^*$ decay, implying S-wave dominated nonfactorizable effects. The nonfactorizable term dominated by the S-wave is also essential for understanding the decay rate of $B^-\to D^{*0}\rho^-$. It is found that all nonfactorizable effects $A_1^{nf}/A_1^{BK^*},\ A_1^{nf}/A_1^{B\rho},\ A_1^{nf}/A_1^{BD^*}$ (nf standing for nonfactorization) are positive and of order 10%, in accordance with $\chi_2(B\to D(D^*)\pi(\rho))$ and $\chi_2(B\to\psi K)$. However, we show that in $D\to\bar{K}^*\rho$ decay nonfactorizable effects cannot be dominated by the S-wave. A polarization measurement in the color- and Cabibbo-suppressed decay mode $D^+\to\phi\rho^+$ is strongly urged in order to test if A_2^{nf}/A_2 plays a more pivotal role than A_1^{nf}/A_1 in charm decay.

1. Introduction

It is customary to assume that two-body nonleptonic weak decays of heavy mesons are dominated by factorizable contributions. Under this assumption, the spectator meson decay amplitude is the product of the universal parameter a_1 (for external W-emission) or a_2 (for internal W-emission), which is channel independent in D or B decays, and hadronic matrix elements which can be factorized as the product of two independent hadronic currents. The universal parameters a_1 and a_2 are related to the Wilson coefficient functions c_1 and c_2 by

$$a_1 = c_1 + \frac{1}{N_c}c_2, \qquad a_2 = c_2 + \frac{1}{N_c}c_1, \tag{1}$$

with N_c being the number of colors. It is known that the bulk of exclusive nonleptonic charm decay data cannot be explained by this factorization approach [1]. For example, the predicted ratio of the color-suppressed mode $D^0 \to \bar{K}^0\pi^0$ and color-favored decay $D^0 \to K^-\pi^+$ is in violent disagreement with experiment. This signals the importance of the nonfactorizable effects.

The leading nonfactorizable contribution arises from the soft gluon exchange between two color-octet currents

$$O_c = \frac{1}{2}(\bar{q}_1\lambda^a q_2)(\bar{q}_3\lambda^a q_4), \tag{2}$$

where $(\bar{q}_1\lambda^a q_2)$ stands for $\bar{q}_1\gamma_\mu(1-\gamma_5)\lambda^a q_2$. For $M \to PP$, VP decays (P: pseudoscalar meson, V: vector meson), the nonfactorizable effect amounts to a redefinition of the parameters a_1 and a_2 [2], [1]

$$a_1 \to c_1 + c_2(\frac{1}{N_c} + \chi_1), \qquad a_2 \to c_2 + c_1(\frac{1}{N_c} + \chi_2), \tag{3}$$

where χ_1 and χ_2 denote the contributions of O_c to color-favored and color-suppressed decay amplitudes respectively relative to the factorizable ones. For example, for $D_s^+ \to \phi\pi^+$, $D^+ \to \phi\pi^+$ decays,

$$\begin{aligned} \chi_1(D_s^+ \to \phi\pi^+) &= \frac{\langle\phi\pi^+|\frac{1}{2}(\bar{u}\lambda^a d)(\bar{s}\lambda^a c)|D_s^+\rangle}{\langle\phi\pi^+|(\bar{u}d)(\bar{s}c)|D_s^+\rangle_f}, \\ \chi_2(D^+ \to \phi\pi^+) &= \frac{\langle\phi\pi^+|\frac{1}{2}(\bar{u}\lambda^a c)(\bar{s}\lambda^a s)|D^+\rangle}{\langle\phi\pi^+|(\bar{u}c)(\bar{s}s)|D^+\rangle_f}. \end{aligned} \tag{4}$$

The subscript f in Eq.(4) denotes a factorizable contribution:

$$\begin{aligned} \langle\phi\pi^+|(\bar{u}d)(\bar{s}c)|D_s^+\rangle_f &= 2m_\phi f_\pi(\varepsilon^*\cdot p_{D_s})A_0^{D_s\phi}(m_\pi^2), \\ \langle\phi\pi^+|(\bar{u}c)(\bar{s}s)|D^+\rangle_f &= m_\phi f_\phi(\varepsilon^*\cdot p_D)F_1^{D\pi}(m_\phi^2), \end{aligned} \tag{5}$$

[1] Note that our definition of χ_1 and χ_2 is different from r_1 and r_2 defined in [3] by a factor of 2.

where ε_μ is the polarization vector of the ϕ meson, and we have followed Ref.[4] for the definition of form factors. The nonfactorizable contributions have the expressions

$$\begin{aligned}\langle\phi\pi^+|\frac{1}{2}(\bar{u}\lambda^a d)(\bar{s}\lambda^a c)|D_s^+\rangle &= 2m_\phi f_\pi(\varepsilon^*\cdot p_{_{D_s}})A_0^{nf}(m_\pi^2),\\ \langle\phi\pi^+|\frac{1}{2}(\bar{u}\lambda^a c)(\bar{s}\lambda^a s)|D^+\rangle &= m_\phi f_\phi(\varepsilon^*\cdot p_{_D})F_1^{nf}(m_\phi^2),\end{aligned}\tag{6}$$

with the superscript nf referring to nonfactorizable contributions. It is clear that

$$\chi_1(D_s^+\to\phi\pi^+)=\frac{A_0^{nf}(m_\pi^2)}{A_0^{D_s\phi}(m_\pi^2)},\qquad \chi_2(D^+\to\phi\pi^+)=\frac{F_1^{nf}(m_\phi^2)}{F_1^{D\pi}(m_\phi^2)}.\tag{7}$$

That is, χ simply measures the fraction of nonfactorizable contributions to the form factor under consideration.

Although we do not know how to calculate χ_1 and χ_2 from first principles, we do anticipate that [3]

$$|\chi(B\to PP)|<|\chi(D\to PP)|<|\chi(D\to VP)|,\tag{8}$$

based on the reason that nonperturbative soft gluon effects become more important when the final-state particles move slower, allowing more time for significant final-state interactions after hadronization. As a consequence, it is obvious that $a_{1,2}$ are in general not universal and that the rule of discarding $1/N_c$ terms [5], which works empirically well in $D\to\bar{K}\pi$ decay, cannot be safely extrapolated to $B\to D\pi$ decay as $|\chi(B\to D\pi)|$ is expected to be much smaller than $|\chi(D\to\bar{K}\pi)\sim-\frac{1}{3}|$ (the c.m. momentum in $D\to\bar{K}\pi$ being 861 MeV, to be compared with 2307 MeV in $B\to D\pi$) and hence a large cancellation between $1/N_c$ and $\chi(B\to D\pi)$ is not expected to happen. The recent CLEO observation [6] that the rule of discarding $1/N_c$ terms is not operative in $B\to D(D^*)\pi(\rho)$ decays is therefore not stunning. Only the fact that $\chi(B\to D\pi)$ is positive turns out to be striking.

Unlike the PP or VP case, it is not pertinent to define $\chi_{1,2}$ for $M\to VV$ decay as its general amplitude consists of three independent Lorentz scalars:

$$A[M(p)\to V_1(\varepsilon_1,p_1)V_2(\varepsilon_2,p_2)]\propto\varepsilon_\mu^*(\lambda_1)\varepsilon_\nu^*(\lambda_2)(\hat{A}_1g^{\mu\nu}+\hat{A}_2p^\mu p^\nu+i\hat{V}\epsilon^{\mu\nu\alpha\beta}p_{1\alpha}p_{2\beta}),\tag{9}$$

where $\hat{A}_1,\ \hat{A}_2,\ \hat{V}$ are related to the form factors $A_1,\ A_2$ and V respectively. Since *a priori* there is no reason to expect that nonfactorizable terms weight in the same way to S-, P- and D-waves, namely $A_1^{nf}/A_1=A_2^{nf}/A_2=V^{nf}/V$, we thus cannot define χ_1 and χ_2. Consequently, it is in general not possible to define an effective a_1 or a_2 for $M\to VV$ decays once nonfactorizable effects are taken into account [7]. In the factorization approach, the fraction of polarization, say Γ_L/Γ (L: longitudinal polarization) in $B\to\psi K^*$ decay, is independent

of the parameter a_2. As a result, if an effective a_2 can be defined for $B \to \psi K^*$, it will lead to the conclusion that nonfactorizable terms cannot affect the factorization prediction of Γ_L/Γ at all. It was realized recently that all the known models in the literature in conjunction with the factorization hypothesis fail to reproduce the data of Γ_L/Γ or the production ratio $\Gamma(B \to \psi K^*)/\Gamma(B \to \psi K)$ or both [8,9]. Evidently, if we wish to utilize nonfactorizable effects to resolve the puzzle with Γ_L/Γ, a key ingredient will be the nonexistence of an effective a_2 for $B \to \psi K^*$.

In short, there are two places where the factorization hypothesis can be unambiguiously tested: (i) To extract the parameters a_1 and a_2 from the experimental measurements of $M \to PP,\ VP$ to see if they are process independent. (ii) To measure the fraction of longitudinal polarization in $M \to VV$ decay and compare with the factorization prediction. Any failure of them will indicate a breakdown of factorization.

The purpose of the present paper is threefold. (i) The parameters χ_1 and χ_2 have been extracted in Ref.[3] (see also [10]). Here we wish to update the values of $\chi_{1,2}$ using the q^2 dependence of form factors suggested by QCD-sum-rule calculations and other theoretical arguments. (ii) It was recently advocated by Kamal and Sandra [7] that the assumption that in $B \to \psi K^*$ decay the nonfactorizable amplitude contributes only to S-wave final states, namely $A_1^{nf} \neq 0,\ A_2^{nf} = V^{nf} = 0$, will lead to a satisfactory explanation of the data of $\Gamma(B \to \psi K^*)/\Gamma(B \to \psi K)$ and Γ_L/Γ. We would like to show that this very assumption is also essential for understanding the ratio $\mathcal{B}(B^- \to D^{*0}\rho^-)/\mathcal{B}(\bar{B}^0 \to D^{*+}\rho^-)$, which cannot be explained satisfactorily in previous work. (iii) Contrary to the B meson case, we will demonstrate that the assumption of S-wave dominated nonfactorizable terms does not work in $D \to VV$ decay.

2. Nonfactorizable contributions in $D,\ B \to PP,\ VP$ decays

Because of the presence of final-state interactions (FSI) and the nonspectator contributions (W-exchange and W-annihilation), it is generally not possible to extract the nonfactorization parameters $\chi_{1,2}$ except for a very few channels. Though color-suppressed decays, for example, $D^0 \to \bar{K}^0(\bar{K}^{*0})\pi^0(\rho^0)$ are conventionally classified as Class II modes [11], color-flavored decay $D^0 \to K^-\pi^+$ will bring some important contribution to $D^0 \to \bar{K}^0\pi^0$ via FSI. This together with the small but not negligible W-exchange amplitude renders the determination of a_2 from $D^0 \to \bar{K}^0\pi^0$ impossible. Therefore, in order to determine a_1 and especially a_2 we should focus on the exotic channels e.g. $D^+ \to \bar{K}^0\pi^+,\ \pi^+\pi^0$, and the decay modes with one single isospin component, e.g. $D^+ \to \pi^+\phi,\ D_s^+ \to \pi^+\phi$, where nonspectator contributions are absent and FSI are presumably negligible.

We next write down the relations between $\chi_{1,2}$ and form factors

$$\begin{aligned}
\chi_1(D\to\bar K\pi) &= \frac{F_0^{nf}(m_\pi^2)}{F_0^{DK}(m_\pi^2)}, &\qquad \chi_2(D\to\bar K\pi) &= \frac{F_0^{nf}(m_K^2)}{F_0^{D\pi}(m_K^2)},\\
\chi_1(D\to\bar K^*\pi) &= \frac{A_0^{nf}(m_\pi^2)}{A_0^{DK^*}(m_\pi^2)}, &\qquad \chi_2(D\to\bar K^*\pi) &= \frac{F_1^{nf}(m_{K^*}^2)}{F_1^{D\pi}(m_{K^*}^2)},\\
\chi_1(D\to\bar K\rho) &= \frac{F_1^{nf}(m_\rho^2)}{F_1^{DK}(m_\rho^2)}, &\qquad \chi_2(D\to\bar K\rho) &= \frac{A_0^{nf}(m_K^2)}{A_0^{D\rho}(m_K^2)},\\
\chi_1(D_s^+\to\phi\pi^+) &= \frac{A_0^{nf}(m_\pi^2)}{A_0^{D_s\phi}(m_\pi^2)}, &\qquad \chi_2(D^+\to\phi\pi^+) &= \frac{F_1^{nf}(m_\phi^2)}{F_1^{D\pi}(m_\phi^2)}.
\end{aligned} \tag{10}$$

It is clear that only the three form factors $F_0,\ F_1$ and A_0 entering into the decay amplitudes of $M\to PP,\ VP$. A consideration of the heavy quark limit behavior of the form factors suggests that the q^2 dependence of F_1 (A_2) is different from that of F_0 (A_0 and A_1) by an additional pole factor [12]. Indeed, QCD-sum-rule calculations have implied a monopole behavior for $F_1(q^2)$ [13-16] and an approximately constant F_0 [15]. With a dipole form factor A_2, as shown by a recent QCD-sum-rule analysis [16], we will thus assume a monopole behavior for A_0.

Unlike the decays $D^+\to\pi^+\phi,\ D_s^+\to\pi^+\phi$ which are described by a single quark diagram, we cannot extract $\chi_{1,2}$ from the data of $D^+\to\bar K^0\pi^+,\ \bar K^0\rho^+,\ \bar K^{*0}\pi^+$ alone without providing further information. For example, the decay amplitude of $D^+\to\bar K^0\pi^+$ reads

$$A(D^+\to\bar K^0\pi^+) = \frac{G_F}{\sqrt2}V_{cs}^*V_{ud}[a_1(m_D^2-m_K^2)f_\pi F_0^{DK}(m_\pi^2)+a_2(m_D^2-m_\pi^2)f_K F_0^{D\pi}(m_K^2)], \tag{11}$$

which consists of external W-emission and internal W-emission amplitudes. We will therefore make a plausible assumption that $\chi_1\sim\chi_2$ so that $\chi(D\to\bar K\pi)$ can be determined from the measured rate of $D^+\to\bar K^0\pi^+$. Since the extraction procedure is already elucidated in Ref.[3], here we will simply present the results (only the central values being quoted) followed by several remarks

$$\begin{aligned}
\chi_2(D\to\bar K\pi) &\simeq -0.36\,,\\
\chi_2(D\to\bar K^*\pi) &\simeq -0.61\,,\\
\chi_2(D^+\to\phi\pi^+) &\simeq -0.44\,,
\end{aligned} \tag{12}$$

where we have used the following quantities:

$$\begin{aligned}
& c_1(m_c)=1.26, \qquad c_2(m_c)=-0.51,\\
& f_\pi=132\,\mathrm{MeV}, \quad f_K=160\,\mathrm{MeV}, \quad f_{K^*}=220\,\mathrm{MeV}, \quad f_\phi=237\,\mathrm{MeV},\\
& F_0^{DK}(0)=F_1^{DK}(0)=0.77\pm0.04\ [17], \qquad F_0^{D\pi}(0)=F_1^{D\pi}(0)=0.83\ [18],\\
& A_1^{DK^*}(0)=0.61\pm0.05, \qquad A_2^{DK^*}(0)=0.45\pm0.09\ [17], \ \Rightarrow A_0^{DK^*}(0)=0.70,
\end{aligned} \tag{13}$$

and the Particle Data Group [19] for the decay rates of various decay modes.

Several remarks are in order. (i) As pointed out by Soares [10], the solutions for χ are not uniquely determined. For example, the other possible solution for $\chi_2(D \to \bar{K}\pi)$ is -1.18. To remove the ambiguities, we have assumed that nonfactorizable corrections are small compared to the factorizable ones. (ii) Assuming $A_0^{D\rho}(0) = A_0^{DK^*}(0)$, we find from the decay $D^+ \to \bar{K}^0\rho^+$ that $\chi(D \to \bar{K}\rho) \approx -1.5$, wich is unreasonably too large. We do not know how to resolve this problem except for noting that thus far there is only one measurement of this decay mode [20]. (iii) To determine $\chi_1(D_s^+ \to \phi\pi^+)$ requires a better knowledge of the form factor $A_0^{D_s\phi}$ and the branching ratio of $D_s^+ \to \phi\pi^+$. Unfortunatly, a direct measurement of them is still not available. Assuming $A_0^{D_s\phi}(0) \approx A_0^{DK^*}(0)$ and $\mathcal{B}(D_s^+ \to \phi\pi^+) = (3.5 \pm 0.4)\%$ [19], we get $\chi_1(D_s^+ \to \phi\pi^+) \approx -0.60$. So in general nonfactorizable terms are process or class dependent, and satisfy the relation $|\chi(D \to PP)| < |\chi(D \to VP)|$ as expected. (iv) Since $\chi_2(D \to \bar{K}\pi)$ is close to $-\frac{1}{3}$, it is evident that a large cancellation between $1/N_c$ and $\chi_2(D \to \bar{K}\pi)$ occurs. This is the dynamic reason why the large-N_c approach operates well for $D \to \bar{K}\pi$ decay. However, this is no longer the case for $D \to VP$ decays. The predicted branching ratios in $1/N_c$ expansion are

$$\begin{aligned} \mathcal{B}(D^+ \to \bar{K}^{*0}\pi^+) &= 0.3\%, \qquad & \mathcal{B}(D^+ \to \bar{K}^0\rho^+) &= 16\%, \\ \mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+) &= 17\%, \qquad & \mathcal{B}(D^+ \to \phi\pi^+) &= 0.4\%, \end{aligned} \tag{14}$$

to be compared with data [19]

$$\begin{aligned} \mathcal{B}(D^+ \to \bar{K}^{*0}\pi^+)_{\rm expt} &= (2.2 \pm 0.4)\%, \qquad & \mathcal{B}(D^+ \to \bar{K}^0\rho^+)_{\rm expt} &= (6.6 \pm 2.5)\%, \\ \mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+)_{\rm expt} &= (4.8 \pm 1.8)\%, \qquad & \mathcal{B}(D^+ \to \phi\pi^+)_{\rm expt} &= (0.67 \pm 0.08)\%. \end{aligned} \tag{15}$$

Consider the decay $D^+ \to \bar{K}^{*0}\pi^+$ as an example. Its amplitude is given by

$$A(D^+ \to \bar{K}^{*0}\pi^+) = \sqrt{2}G_F V_{cs}^* V_{ud}[a_1 f_\pi m_{K^*} A_0^{DK^*}(m_\pi^2) + a_2 f_{K^*} m_{K^*} F_1^{D\pi}(m_{K^*}^2)]. \tag{16}$$

Since the interference is destructive and $f_{K^*}F_1^{D\pi} > f_\pi A_0^{DK^*}$, a large $|a_2|$ is needed in order to enhance the branching ratio of $D^+ \to \bar{K}^{*0}\pi^+$ from 0.3% to 2.2%. (Note that a_1 is relatively insensitive to the nonfactorizable effects.) This in turn implies a negative $(\frac{1}{N_c} + \chi_2)$ and hence $\chi_2(D \to \bar{K}^*\pi) < -\frac{1}{3}$. Therefore, we are led to conclude that the leading $1/N_c$ expansion cannot be a universal approach for the nonleptonic weak decays of the meson. However, the fact that substantial nonfactorizable effects which contribute destructively with the subleading $1/N_c$ factorizable contributions are required to accommodate the data of charm decay means that, as far as charm decays are concerned, the large-N_c approach greatly improves the naive factorization method in which $\chi_{1,2} = 0$; the former approach amounts to having a universal nonfactorizable term $\chi_{1,2} = -1/N_c$.

We next turn to $B \to D(D^*)\pi(\rho)$ decays. Though both nonspectator and FSI effects are known to be important in charm decays, it is generally believed that they do not play a significant role in bottom decays as the decay particles are moving fast, not allowing adequate time for FSI. This gives the enormous advantage that it is conceivable to determine a_1 and a_2 separately from $B \to D(D^*)\pi(\rho)$ decays. Using the heavy-flavor-symmetry approach for heavy-light form factors and assuming a monopole extrapolation for $F_1,\ A_0,\ A_1$, a dipole behavior for $A_2,\ V$, and an approximately constant F_0, as suggested by QCD-sum-rule calculations and some theoretical arguments [21], we found from the CLEO data that [21] [2]

$$\begin{aligned} a_1(B \to D^{(*)}\pi(\rho)) &= 1.01 \pm 0.06\,, \\ a_2(B \to D^{(*)}\pi(\rho)) &= 0.23 \pm 0.06\,. \end{aligned} \tag{17}$$

Taking $c_1(m_b) = 1.11$ and $c_2(m_b) = -0.26$ leads to

$$\chi_1(B \to D^{(*)}\pi(\rho)) \simeq 0.05\,, \qquad \chi_2(B \to D^{(*)}\pi(\rho)) \simeq 0.11\,. \tag{18}$$

Since $(\frac{1}{N_c} + \chi_{1,2}) = (a_{1,2} - c_{1,2})/c_{2,1}$ and $|c_2| << |c_1|$, it is clear that the determination of χ_1 is far more uncertain than χ_2: it is very sensitive to the values of $a_1,\ c_1$ and c_2. We see from (18) that nonfactorizable effects become less important in B decays, as what expected [see (8)]. However, a positive $\chi_2(B \to D(D^*)\pi(\rho))$, which is necessary to explain the constructive interference in $B^- \to D^0(D^{*0})\pi^-(\rho^-)$ decays, appears to be rather striking. A recent light cone QCD-sum-rule calculation [22] following the framework outlined in [23] fails to reproduce a positive $\chi_2(B \to D\pi)$. This tantalizing issue should be resolved in the near future.

For $B \to \psi K$ decays, we found [21]

$$\Big|a_2(B^- \to \psi K^-)\Big| = 0.235 \pm 0.018\,, \qquad \Big|a_2(B^0 \to \psi K^0)\Big| = 0.192 \pm 0.032\,. \tag{19}$$

The combined value is

$$a_2(B \to \psi K) = 0.225 \pm 0.016\,, \tag{20}$$

where its sign should be positive, as we have argued in [21]. (It was advocated by Soares [10] that an analysis of the contribution of $B \to \psi K$ to the decay $B \to K\ell^+\ell^-$ can be used to remove the sign ambiguity of a_2.) It follows that

$$\chi_2(B \to \psi K) = \frac{F_1^{nf}(m_\psi^2)}{F_1^{BK}(m_\psi^2)} \simeq 0.10\,, \tag{21}$$

[2]Contrary to the charmed meson case, the variation of $a_{1,2}$ from $B \to D\pi$ to $D^*\pi$ and $D\rho$ decays is negligible (see Table IV of [21]).

which is in accordance with $\chi_2(B \to D^{(*)}\pi(\rho))$.

Finally, it is very interesting to note that, in contrast to charm decays, the large-N_c approach is even worse than the naive factorization method in describing $B \to D(D^*)\pi(\rho)$ decays as $\chi_2(B \to D^{(*)}\pi(\rho))$ is small but positive.

3. Nonfactorizable contributions in $B \to \psi K^*,\ D^*\rho$ decays

As stressed in the Introduction, in general one cannot define $\chi_{1,2}$ and hence an effective $a_{1,2}$ for $M \to VV$ decays unless the nonfactorizable terms weight in the same manner in all three partial waves. It was pointed out recently that there are two experimental data, namely the production ratio $R \equiv \Gamma(B \to \psi K^*)/\Gamma(B \to \psi K)$ and the fraction of longitudinal polarization Γ_L/Γ in $B \to \psi K^*$, which cannot be accounted for simultaneously by all commonly used models within the framework of factorization [8,9]. The experimental results are

$$R = 1.74 \pm 0.39 \text{ [6]}, \qquad \frac{\Gamma_L}{\Gamma} = 0.78 \pm 0.07\,, \tag{22}$$

where the latter is the combined average of the three measurements:

$$\left(\frac{\Gamma_L}{\Gamma}\right)_{B\to\psi K^*} = \begin{cases} 0.97 \pm 0.16 \pm 0.15, & \text{ARGUS [24];} \\ 0.80 \pm 0.08 \pm 0.05, & \text{CLEO [6];} \\ 0.66 \pm 0.10^{+0.08}_{-0.10}, & \text{CDF [25].} \end{cases} \tag{23}$$

Irrespective of the production ratio R, all the existing models fail to produce a large longitudinal polarization fraction [8,9]. This strongly implies that the puzzle with Γ_L/Γ can only be resolved by appealing to nonfactorizable effects. [3] However, if the relation $A_1^{nf}/A_1 = A_2^{nf}/A_2 = V^{nf}/V$ holds, then an effective a_2 can be defined for $B \to \psi K^*$ and the prediction of Γ_L/Γ will be the same as that in the factorization approach as the polarization fraction is independent of a_2. Consequently, nonfactorizable terms should contribute differently to S-, P- and D-wave amplitudes if we wish to explain the observed Γ_L/Γ.

The large longitudinal polarization fraction observed by ARGUS and CLEO suggests that the decay $B \to \psi K^*$ is almost all S-wave. To see this, we write down the $B \to \psi K^*$

[3] An interesting observation was made recently in [26] that the factorization assumption in $B \to \psi K(K^*)$ is not ruled out and the data can be accommodated by the heavy-flavor-symmetry approach for heavy-light form factors provided that the $A_1(q^2)$ form factor is frankly decreasing. To our knowledge, a decreasing A_1 with q^2 is ruled out by several recent QCD-sum-rule analyses (see e.g. [16]). Using the same approach for heavy-light form factors but the q^2 dependence of form factors given in [21], we found that $R = 1.84$ and $\Gamma_L/\Gamma = 0.56$ [21]. Evidently, the factorization approach is still difficult to explain the observed large polarization fraction.

amplitude

$$
\begin{aligned}
A[B(p)\to\psi(p_1)K^*(p_2)] &= \frac{G_F}{\sqrt{2}}V_{cs}^*V_{ud}\left(c_2+\frac{c_1}{3}\right)f_\psi m_\psi\varepsilon_\mu^*(\psi)\varepsilon_\nu^*(K^*)[\hat{A}_1 g^{\mu\nu}+\hat{A}_2 p^\mu p^\nu \\
&+ i\hat{V}\epsilon^{\mu\nu\alpha\beta}p_{1\alpha}p_{2\beta}], \qquad (24)
\end{aligned}
$$

with

$$
\begin{aligned}
\hat{A}_1 &= (m_B+m_{K^*})A_1^{BK^*}(m_\psi^2)\left[1+\kappa\frac{A_1^{nf}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)}\right], \\
\hat{A}_2 &= -\frac{2}{(m_B+m_{K^*})}A_2^{BK^*}(m_\psi^2)\left[1+\kappa\frac{A_2^{nf}(m_\psi^2)}{A_2^{BK^*}(m_\psi^2)}\right], \qquad (25) \\
\hat{V} &= -\frac{2}{(m_B+m_{K^*})}V^{BK^*}(m_\psi^2)\left[1+\kappa\frac{V^{nf}(m_\psi^2)}{V^{BK^*}(m_\psi^2)}\right],
\end{aligned}
$$

and $\kappa = c_1/(c_2+\frac{1}{3}c_1)$. It is easily seen that we will have an effective $a_2 = c_2 + c_1(\frac{1}{3}+\chi_2)$ if the nonfactorizable terms happen to satisfy the relation $A_1^{nf}/A_1 = A_2^{nf}/A_2 = V^{nf}/V = \chi_2$. The decay rate of this mode is of the form

$$
\Gamma(B\to\psi K^*) \propto\ (a-b\tilde{x})^2 + 2(1+c^2\tilde{y}^2), \qquad (26)
$$

where

$$
\begin{aligned}
a &= \frac{m_B^2-m_\psi^2-m_{K^*}^2}{2m_\psi m_{K^*}}, \qquad b = \frac{2m_B^2p_c^2}{m_\psi m_{K^*}(m_B+m_{K^*})^2}, \qquad c = \frac{2m_Bp_c}{(m_B+m_{K^*})^2}, \\
\tilde{x} &= \frac{A_2^{BK^*}(m_\psi^2)+\kappa A_2^{nf}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)+\kappa A_1^{nf}(m_\psi^2)}, \qquad \tilde{y} = \frac{V^{BK^*}(m_\psi^2)+\kappa V^{nf}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)+\kappa A_1^{nf}(m_\psi^2)}, \qquad (27)
\end{aligned}
$$

with p_c being the c.m. momentum. The longitudinal polarization fraction is then given by

$$
\frac{\Gamma_L}{\Gamma} = \frac{(a-b\tilde{x})^2}{(a-b\tilde{x})^2+2(1+c^2\tilde{y}^2)}. \qquad (28)
$$

If the decay is an almost S-wave, one will have $\Gamma_L/\Gamma \sim a^2/(a^2+2) = 0.83$. Since $\kappa >> 1$, $\tilde{x}$ (D-wave) and $\tilde{y}$ (P-wave) can be suppressed by assuming that, as first postulated in [7], in $B\to\psi K^*$ decay the nonfactorizable amplitude contributes only to S-wave final states; that

is, [4]

$$A_1^{nf} \neq 0, \qquad A_2^{nf} = V^{nf} = 0. \tag{29}$$

The rational for this assumption is given in [7].

With the assumption (29), the branching ratio followed from (24) is

$$\mathcal{B}(B \to \psi K^*) = 0.0288 \left| \left(c_2 + \frac{c_1}{3} \right) A_1^{BK^*}(m_\psi^2) \right|^2 [(a\xi - bx)^2 + 2(\xi^2 + c^2 y^2)] \tag{30}$$

with

$$x = \frac{A_2^{BK^*}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)}, \quad y = \frac{V^{BK^*}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)}, \quad \xi = 1 + \kappa \frac{A_1^{nf}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)}, \tag{31}$$

where uses of $|V_{cb}| = 0.040$ and $\tau(B) = 1.52 \times 10^{-12} s$ have been made. It follows that

$$\frac{\Gamma_L}{\Gamma} = \frac{(a\xi - bx)^2}{(a\xi - bx)^2 + 2(1 + c^2 y^2)}. \tag{32}$$

We use the measured branching ratio $\mathcal{B}(B \to \psi K^*) = (0.172 \pm 0.030)\%$ [6] to determine the ratio $A_1^{nf}(m_\psi^2)/A_1^{BK^*}(m_\psi^2)$, which is found to be

$$\frac{A_1^{nf}(m_\psi^2)}{A_1^{BK^*}(m_\psi^2)} \simeq 0.08\,, \tag{33}$$

which we have used $A_1^{BK^*}(m_\psi^2) = 0.41$, $A_2^{BK^*}(m_\psi^2) = 0.36$, $V^{BK^*}(m_\psi^2) = 0.72$ [21] and discarded the other possible solution $A_1^{nf}/A_1^{BK^*} = -0.22$ for its "wrong" sign, recalling that F_1^{nf}/F_1^{BK} is positive [cf. Eq.(21)]. The predicted longitudinal polarization fraction is $\Gamma_L/\Gamma = 0.73$, which is in accordance with experiment.

The assumption of negligible nonfactorizable contributions to P- and D-waves also turns out to be essential for understanding the decay rate of $B^- \to D^{*0}\rho^-$ or the ratio $R_4 \equiv \mathcal{B}(B^- \to D^{*0}\rho^-)/\mathcal{B}(\bar{B}^0 \to D^{*+}\rho^-)$. The issue arises as follows. In Ref.[21] we have determined a_1 and a_2 from $B \to D\pi,\ D^*\pi,\ D\rho$ decays and obtained a consistent ratio a_2/a_1 from

[4] A different approach for nonfactorizable effects adopted in Ref.[27] amounts to $A_1^{nf} = A_2^{nf} = 0$ and $V^{nf} \neq 0$. It follows from Eq.(28) that

$$\frac{\Gamma_L}{\Gamma} = \frac{(a - bx)^2}{(a - bx)^2 + 2(1 + c^2\bar{y}^2)},$$

with $\bar{y} = (V^{BK^*}(m_\psi^2) + \kappa V^{nf}(m_\psi^2))/A_1^{BK^*}(m_\psi^2)$ and x being defined in (31). It is clear that in order to get a large longitudinal polarization fraction one needs a *negative* V^{nf}/V ! Using the numerical values $a = 3.164,\ b = 1.304,\ x = 0.89$, we find $(\Gamma_L/\Gamma)_{\max} = 0.67$. The prediction $\Gamma_L/\Gamma = 0.65$ given by [27] is one standard deviation from experiment (22).

channel to channel: $0.24 \pm 0.10,\ 0.24 \pm 0.14,\ 0.21 \pm 0.08$ (see Table IV of [21]). Assuming factorization, we got $a_2/a_1 = 0.34 \pm 0.13$ from $B \to D^*\rho$ decay, which deviates somewhat from above values. In the presence of S-wave dominated nonfactorizable contributions, it is no longer possible to define an effective a_1 and a_2 for $B \to D^*\rho$ decay. Therefore, the quantities to be compared with are A_1^{nf}/A_1 in $B \to D^*\rho$ decay and χ_2 in $B \to D\pi,\ D^*\pi,\ D\rho$. A straightforward calculation yields (see [21] for the factorizable case)

$$R_4 = \frac{\tau(B^-)}{\tau(B^0)} \left(1 + 2\eta \frac{H_1}{H} + \eta^2 \frac{H_2}{H}\right), \tag{34}$$

with

$$\begin{aligned} H &= (\hat{a}\hat{\xi} - \hat{b}\hat{x})^2 + 2(\hat{\xi}^2 + \hat{c}^2\hat{y}^2), \\ H_1 &= (\hat{a}\hat{\xi} - \hat{b}\hat{x})(\hat{a}\hat{\xi}' - \hat{b}'\hat{x}') + 2(\hat{\xi}\hat{\xi}' + \hat{c}\hat{c}'\hat{y}\hat{y}'), \\ H_2 &= (\hat{a}\hat{\xi}' - \hat{b}'\hat{x}')^2 + 2(\hat{\xi}'^2 + \hat{c}'^2\hat{y}'^2), \\ \eta &= \frac{m_{D^*}(m_B + m_\rho)}{m_\rho(m_B + m_{D^*})} \frac{f_{D^*}}{f_\rho} \frac{A_1^{B\rho}(m_{D^*}^2)}{A_1^{BD^*}(m_\rho^2)} \frac{c_2 + \frac{1}{3}c_1}{c_1 + \frac{1}{3}c_2}, \\ \hat{\xi} &= 1 + \frac{c_2}{c_1 + \frac{1}{3}c_2} \frac{A_1^{nf}(m_\rho^2)}{A_1^{BD^*}(m_\rho^2)}, \\ \hat{\xi}' &= 1 + \frac{c_1}{c_2 + \frac{1}{3}c_1} \frac{A_1^{nf}(m_{D^*}^2)}{A_1^{B\rho}(m_{D^*}^2)}, \end{aligned} \tag{35}$$

where $\hat{a},\ \hat{b},\ \hat{c}$ are obtained from $a,\ b,\ c$ respectively in (27), $\hat{x},\ \hat{y}$ from $x,\ y$ in (31) by replacing $\psi \to D^*,\ K^* \to \rho$, and $\hat{b}',\ \hat{c}',\ \hat{x}',\ \hat{y}'$ are obtained from $\hat{b},\ \hat{c},\ \hat{x},\ \hat{y}$ respectively by replacing $D^* \leftrightarrow \rho$; for instance $\hat{x}' = A_2^{B\rho}(m_{D^*}^2)/A_1^{B\rho}(m_{D^*}^2)$. Assuming $A_1^{nf}/A_1^{BD^*} \sim A_1^{nf}/A_1^{B\rho}$ and fitting (34) to the experimental value $R_4 = (1.68 \pm 0.35)\%$ [6], we get

$$\frac{A_1^{nf}(m_{D^*}^2)}{A_1^{B\rho}(m_{D^*}^2)} \sim \frac{A_1^{nf}(m_\rho^2)}{A_1^{BD^*}(m_\rho^2)} \simeq 0.12\,. \tag{36}$$

We see that the S-wave dominated nonfactorizable effect in $B \to \psi K^*$ and $B \to D^*\rho$ decays is of order 10%, consistent with $\chi_2(B \to \psi K)$ and $\chi_2(B \to D(D^*)\pi(\rho))$.

4. Nonfactorizable contributions in $D \to \bar{K}^*\rho$ decay

We have shown in the previous section that S-wave dominated nonfactorizable terms are needed to explain the large longitudinal polarization fraction observed in $B \to \psi K^*$ and the ratio $\mathcal{B}(B^- \to D^{*0}\rho^-)/\mathcal{B}(\bar{B}^0 \to D^{*+}\rho^-)$. However, we shall see in this section that the assumption (29) is no longer applicable to $D \to \bar{K}^*\rho$ decay. An experimental measurement of $D^+ \to \bar{K}^{*0}\rho^+$ and $D^0 \to \bar{K}^{*0}\rho^0$ by Mark III [28] shows that (i) the decay $D^+ \to \bar{K}^{*0}\rho^+$ is a mixture of longitudinal and transverse polarization consistent with a pure

S-wave amplitude, [5] and (ii) $D^0 \to \bar{K}^{*0}\rho^0$ is almost all transverse, requiring a cancellation between the longitudinal S-wave and D-wave.

We first consider the decay $D^+ \to \bar{K}^{*0}\rho^+$, whose amplitude is given by

$$A(D^+(p) \to \bar{K}^{*0}(p_1)\rho^+(p_2)) = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud} \varepsilon_\mu^*(K^*)\varepsilon_\nu^*(\rho)[\tilde{A}_1 g^{\mu\nu} + \tilde{A}_2 p^\mu p^\nu + i\tilde{V}\epsilon^{\mu\nu\alpha\beta}p_{1\alpha}p_{2\beta}], \quad (38)$$

where

$$\begin{aligned} \tilde{A}_1 &= \left(c_1 + \frac{c_2}{3}\right) f_\rho m_\rho (m_D + m_{K^*}) \left(1 + \frac{c_2}{c_1 + \frac{1}{3}c_2} \frac{A_1^{nf}(m_\rho^2)}{A_1^{DK^*}(m_\rho^2)}\right) A_1^{DK^*}(m_\rho^2) \\ &+ \left(c_2 + \frac{c_1}{3}\right) f_{K^*} m_{K^*} (m_D + m_\rho) \left(1 + \frac{c_1}{c_2 + \frac{1}{3}c_1} \frac{A_1^{nf}(m_{K^*}^2)}{A_1^{D\rho}(m_{K^*}^2)}\right) A_1^{D\rho}(m_{K^*}^2), \end{aligned} \quad (39)$$

and $\tilde{A}_2$ ($\tilde{V}$) is obtained from $\tilde{A}_1$ with the replacements $A_1 \to A_2$ ($A_1 \to V$), $(m_D + m_{K^*}) \to -2/(m_D + m_{K^*})$ and $(m_D + m_\rho) \to -2/(m_D + m_\rho)$. Since $A_1^{nf}/A_1^{DK^*}$ and $A_1^{nf}/A_1^{D\rho}$ are expected to be negative [see Eq.(12)], it is obvious that if nonfactorizable terms are dominated by the S-wave, it will imply a more severe destructive interference in the S-wave amplitude than in P- and D-wave amplitudes, in contradiction to the observation that this decay is almost all S-wave. The branching ratio is calculated to be

$$\mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+) = 0.10 \left|\left(c_1 + \frac{1}{3}c_2\right) A_1^{DK^*}(m_\rho^2)\right|^2 (H' + 2\eta' H_1' + \eta'^2 H_2'), \quad (40)$$

with the expressions of η', H', $H'_{1,2}$ analogous to η, H, $H_{1,2}$ in (35). A fit of (40) to the Mark III data for the branching ratio (37) gives rise to (assuming $A_1^{nf}/A_1^{DK^*} \sim A_1^{nf}/A_1^{D\rho}$)

$$\frac{A_1^{nf}(m_\rho^2)}{A_1^{DK^*}(m_\rho^2)} \sim \frac{A_1^{nf}(m_{K^*}^2)}{A_1^{D\rho}(m_{K^*}^2)} \approx -0.98\,, \quad (41)$$

which is uncomfortably too large. [6] Moreover, the P-wave branching ratio is predicted to be 2.0×10^{-2}, in disagreement with experiment [28]

$$\mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+)_{P-\mathrm{wave}} < 0.5 \times 10^{-2}. \quad (42)$$

It thus appears to us that an almost S-wave $D^+ \to \bar{K}^{*0}\rho^+$ implies that

$$\left|\frac{A_2^{nf}}{A_2^{DK^*(\rho)}}\right|, \quad \left|\frac{V^{nf}}{V^{DK^*(\rho)}}\right| \gtrsim \left|\frac{A_1^{nf}}{A_1^{DK^*(\rho)}}\right|. \quad (43)$$

---

[5] The other measurement by E691 [29] disagrees severely with Mark III on the branching ratio

$$\mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+) = \begin{cases} (4.8 \pm 1.2 \pm 1.4)\%, & \text{Mark III [28];} \\ (2.3 \pm 1.2 \pm 0.9)\%, & \text{E691 [29].} \end{cases} \quad (37)$$

Recall that model calculations tend to give a very large branching ratio of 17% [see Eq.(14)].

[6] A fit to the E691 measurement (37) for the branching ratio yields an even larger value: $A_1^{nf}/A_1^{DK^*} \sim A_1^{nf}/A_1^{D\rho} \approx -1.41$.

Taking $A_1^{nf}/A_1 = A_2^{nf}/A_2 = V^{nf}/V = \chi(D \to \bar{K}^*\rho)$ as an illustration, we obtain

$$\chi(D \to \bar{K}^*\rho) \approx -0.65 \tag{44}$$

and $\mathcal{B}(D^+ \to \bar{K}^{*0}\rho^+)_{P-\rm wave} = 2.0 \times 10^{-3}$, which are certainly more plausible than before.

Another indication for the failure of the S-wave dominated hypothesis for nonfactorizable effects comes from the decay $D^0 \to \bar{K}^{*0}\rho^0$, where $\bar{K}^{*0}$ and ρ^0 are completely transversely polarized, implying a large D-wave which is compensated by the longitudinal S-wave. Recall that the factorizable $D \to VV$ amplitudes have the sailent feature :

$$|S\text{−wave amplitude}| > |P\text{−wave amplitude}| > |D\text{−wave amplitude}|. \tag{45}$$

Since the color-suppressed D-wave amplitude of $D^0 \to \bar{K}^{*0}\rho^0$ is proportional to $[1 + c_1/(c_2 + \frac{1}{3}c_1)(A_2^{nf}/A_2^{D\rho})]$, a large D-wave thus indicates a negative A_2^{nf}/A_2 and

$$\left|\frac{A_1^{nf}}{A_1}\right| << \left|\frac{A_2^{nf}}{A_2}\right|, \qquad \text{or} \;\; \frac{A_1^{nf}}{A_1} \approx 0, \;\;\; \frac{A_2^{nf}}{A_2} \neq 0. \tag{46}$$

Therefore, we see that nonfactorizable terms in charm decay are consistently to be negative [cf. Eqs.(12) and (44)]. Unfortunately, at this point we cannot make a further quantitative analysis due to unknown final-state interactions and W-exchange contributions. A measurement of helicities in $D^0 \to \bar{K}^{*0}\rho^0,\ D^+ \to \phi\rho^+$ will be greatly helpful to pin down the issue. In particular, the color- and Cabibbo-suppressed mode $D^+ \to \phi\rho^+$ is very ideal for this purpose since it is not subject to FSI and nonspectator effects. A polarization measurement in this decay is thus strongly urged (though difficult) in order to test if A_2^{nf}/A_2 plays a more essential role than A_1^{nf}/A_1 in charm decay.

5. Discussion and conclusion

The factorization assumption for hadronic weak decays of mesons can be tested on two different grounds: (i) to extract the effective parameters a_1 and especially a_2 from $M \to PP$, VP decays to see if they are process independent, and (ii) to measure helicities in $M \to VV$ decay. Using the q^2 dependence of form factors suggested by QCD-sum-rule calculations and by some theoretical arguments, we have updated our previous work. It is found that a_2 is evidently not universal in charm decay. The parameter χ_2, which measures the nonfactorizable soft-gluon effect on the color-suppressed deacy amplitude relative to the factorizable one, ranges from $-\frac{1}{3}$ to -0.60 in the decays from $D \to \bar{K}\pi$ to $D^+ \to \phi\pi^+,\ D \to \bar{K}^*\pi$. By contrast, the variation of a_2 in $B \to \psi K,\ B \to D(D^*)\pi(\rho)$ is negligible and nonfactorizable terms $\chi_2(B \to \psi K),\ \chi_2(B \to D^{(*)}\pi(\rho))$ are of order 10% with a positive sign. The pattern for the relative magnitudes of nonfactorizable effects

$$|\chi(B \to PP, VP)| < |\chi(D \to PP)| < |\chi(D \to VP)|$$

is thus well established. This means that nonperturbative soft gluon effects become more important when the final states are less energetic, allowing more time for final-state interactions. This explains why a_2 is class (PP or VP mode) dependent in charm decay, whereas it stays fairly stable in B decay.

Taking factorization as a benchmark, we see that the nonfactorizable terms necessary for describing nonleptonic D and B decays are in opposite directions from the factorization framework. On the one hand, the leading $1/N_c$ expansion, which amounts to a universal $\chi = -\frac{1}{3}$, improves the naive factorization method for charm decays. On the other hand, the naive factorization hypothesis works better than the large-N_c assumption for B decays because nonfactorizable effects are small, being of order 10%. The fact that χ is positive makes it even more clear that the large-N_c approach cannot be extrapolated from D to B physics. Theoretically, the next important task for us is to understand why χ is negative in D decay, while it becomes positive in B decay.

As for $M \to VV$ decay, *a priori* effective $a_{1,2}$ cannot be defined since, as pointed out by Kamal and Sandra, its amplitude (factorizable and nonfactorizable) involves three independent Lorentz scalars, corresponding to S, P and D waves. This turns out to be a nice trade-off for solving the puzzle with the large longitudinal polarization fraction Γ_L/Γ observed in $B \to \psi K^*$, which cannot be accounted for by the factorization hypothesis or by nonfactorizable effects weighted in the same way in all three partial waves, namely $A_1^{nf}/A_1 = A_2^{nf}/A_2 = V^{nf}/V$. A large Γ_L/Γ can be achieved if $B \to \psi K^*$ is almost all S-wave, implying that nonfactorizable contributions are dominated by the S-wave. The same assumption is also needed for understanding the ratio $\mathcal{B}(B^- \to D^{*0}\rho^-)/\mathcal{B}(\bar{B}^0 \to D^{*+}\rho^-)$. We found that all nonfactorizable terms $A_1^{nf}/A_1^{BK^*}$, $A_1^{nf}/A_1^{B\rho}$, $A_1^{nf}/A_1^{BD^*}$ are of order 10% consistent with $\chi_2(B \to D(D^*)\pi(\rho))$ and $\chi_2(B \to \psi K)$.

Surprisingly, the assumption of S-wave dominated nonfactorizable effects is not operative in $D \to \bar{K}^*\rho$ decay, which exhibits again another disparity between B and D physics. We found that A_2^{nf}/A_2 should play a more pivotal role than A_1^{nf}/A_1 in charm decay. We thus urge experimentalists to measure helicities in the color- and Cabibbo-suppressed decay mode $D^+ \to \phi\rho^+$ decay to gain insight in the nonfactorizable effects in $D \to VV$ decay.

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